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Measuring New-Physics Parameters in B Penguin Decays

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Abstract

We examine new-physics (NP) effects in B decays with large $\bar{b} \to \bar{s}$ penguin amplitudes. Decays involving $\bar{b} \to \bar{d}$ penguins are assumed to be unaffected. We consider a model-independent parametrization of such NP. We argue that NP strong phases are negligible relative to those of the standard model. This allows us to describe the NP effects in terms of a small number of effective amplitudes $\mathcal{A}^q_{NP}$ ($q = u, d, s, c$) and corresponding weak phases Φ_q. We then consider pairs of neutral B decays which are related by flavour SU(3) in the standard model. One receives a large $\bar{b} \to \bar{s}$ penguin component and has a NP contribution; the other has a $\bar{b} \to \bar{d}$ penguin amplitude and is unaffected by NP. The time-dependent measurement of these two decays allows the *measurement* of the NP parameters $\mathcal{A}^q_{NP}$ and Φ_q. The knowledge of these parameters allows us to rule out many NP models and thus partially identify the new physics.

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The B-factories BaBar and Belle have already made a large number of measurements involving B decays, and this will continue for a number of years. The principal aim of this activity is to test whether the standard model (SM) explanation of CP violation — a complex phase in the Cabibbo-Kobayashi-Maskawa (CKM) matrix [1] — is correct. This is done by measuring CP violation in the B system in many different processes [2]. Hopefully a discrepancy will be found, giving us the first indication of physics beyond the SM.

New-physics (NP) effects in B decays are necessarily virtual processes. As a result, it is generally assumed that, while B-factories can detect the presence of NP, its identification can only be made at future high-energy colliders, in which the new particles are produced directly. The main purpose of this paper is to show that this is not entirely true. Here we will describe a technique which allows us not only to detect the NP, but also to *measure* its amplitude and phase. This will be an important first step in identifying the new physics, even before it has been seen directly at high-energy colliders.

Recently, there have been several hints of such new physics. First, within the SM, the CP-violating asymmetries in $B_d^0(t) \to J/\psi K_S$ and $B_d^0(t) \to \phi K_S$ are both expected to measure the same quantity $\sin 2\beta$ [3]. However, the Belle measurement of $\sin 2\beta$ in $B_d^0(t) \to \phi K_S$ disagrees with that found in $B_d^0(t) \to J/\psi K_S$ by 3.5σ (there is no discrepancy in the BaBar result) [4]. Indeed, the value of $\sin 2\beta$ extracted from all $\bar{b} \to \bar{s}$ penguin decays is 3.1σ below that from charmonium decays. Second, the various $B \to K\pi$ branching ratios have been measured. If one neglects exchange- and annihilation-type amplitudes, which are expected to be small, within the SM one has $R_c = R_n$ [5], where

$$R_c \equiv \frac{2\bar{\Gamma}(B^+ \to K^+\pi^0)}{\bar{\Gamma}(B^+ \to K^0\pi^+)} \quad , \qquad R_n \equiv \frac{\bar{\Gamma}(B_d^0 \to K^+\pi^-)}{2\bar{\Gamma}(B_d^0 \to K^0\pi^0)} \ . \tag{1}$$

However, current measurements yield [4]

$$R_c = 1.42 \pm 0.18 \quad , \qquad R_n = 0.89 \pm 0.13 \ , \tag{2}$$

yielding a discrepancy of 2.4σ between R_c and R_n. Finally, within the SM all CP-violating triple-product correlations (TP's) in $B \to V_1 V_2$ decays (V_1 and V_2 are vector mesons) are expected to vanish or be very small [6]. However, BaBar sees a TP signal in $B \to \phi K^*$ at 1.7σ [7].

While the above new-physics signals are not yet convincing, they do suggest that NP might be playing a role in these decays. In addition, in all cases, the decays in question ($B \to \phi K^{(*)}$ and $B \to K\pi$) receive significant contributions from $\bar{b} \to \bar{s}$ penguin amplitudes. On the other hand, to date there are no NP signals in processes which receive sizeable contributions from $\bar{b} \to \bar{d}$ penguin amplitudes (e.g. $B_d^0 \to \pi\pi$). In this paper, we therefore make the assumption that NP contributes significantly only to those decays which have large $\bar{b} \to \bar{s}$ penguin amplitudes; decays involving $\bar{b} \to \bar{d}$ penguins are not affected.

Up to now, theoretical work has focussed principally on finding signals of new physics in $\bar{b} \to \bar{s}$ transitions – in fact, there are many such signals. However, if NP is found, we will want to identify it. This requires the determination of the NP parameters. Unfortunately, most NP signals simply indicate that physics beyond the SM is present, but do not allow us to extract its parameters. (In some cases, it is possible to put bounds on the NP parameters [8].) The advantage of the technique described in this paper is that it allows us to *measure* the amplitude and phase of the NP.

Assuming that the new physics affects only the $\bar{b} \to \bar{s}$ penguin amplitudes, the first step is a model-independent parametrization of this NP. We assume that a NP piece is added to the effective Hamiltonian:

$$H_{eff} = H_{SM} + H_{NP} \ , \tag{3}$$

where H_{SM} is the SM effective Hamiltonian [9]. H_{NP} contains four-quark operators with all possible Dirac and colour structures, with the proviso that only $\bar{b} \to \bar{s}$ penguin transitions are affected. That is, the general structure of the operators in H_{NP} is $O_{NP} \sim \bar{s}b\,\bar{q}q$ $(q = u, d, s, c)$, where Lorentz and colour structures have been suppressed. We also assume that the contribution from O_{NP} to any B decay is at most of the same order as the SM penguin amplitude.

Taking into account the two different colour assignments, as well as all possible Lorentz structures, there are a total of 20 dimension-six new-physics operators which contribute to each of the $\bar{b} \to \bar{s}q\bar{q}$ $(q = u, d, s, c)$ transitions [10]. These operators, which can contribute to both tree and penguin amplitudes, can be written as

$$\begin{aligned} \mathcal{H}_{NP}^{q} = \sum_{A,B=L,R} \frac{4G_F}{\sqrt{2}} \, \Big\{ & f_{q,1}^{AB}\, \bar{s}_\alpha \gamma_A b_\beta \, \bar{q}_\beta \gamma_B q_\alpha + f_{q,2}^{AB}\, \bar{s} \gamma_A b \, \bar{q} \gamma_B q \\ & + g_{q,1}^{AB}\, \bar{s}_\alpha \gamma^\mu \gamma_A b_\beta \, \bar{q}_\beta \gamma_\mu \gamma_B q_\alpha + g_{q,2}^{AB}\, \bar{s} \gamma^\mu \gamma_A b \, \bar{q} \gamma_\mu \gamma_B q \\ & + h_{q,1}^{AB}\, \bar{s}_\alpha \sigma^{\mu\nu} \gamma_A b_\beta \, \bar{q}_\beta \sigma_{\mu\nu} \gamma_B q_\alpha + h_{q,2}^{AB}\, \bar{s} \sigma^{\mu\nu} \gamma_A b \, \bar{q} \sigma_{\mu\nu} \gamma_B q \Big\} \ , \end{aligned} \tag{4}$$

where we have defined $\gamma_{R(L)} = \frac{1}{2}(1 \pm \gamma_5)$. Although we have written the tensor operators in the same compact form as the other operators, it should be noted that those with $\gamma_A \neq \gamma_B$ are identically zero. Thus, one can effectively set $h_{q,i}^{LR} = h_{q,i}^{RL} = 0$.

In general, all coefficients in Eq. (4) can have new CP-violating weak phases and the matrix elements of the operators will have (process-dependent) CP-conserving strong phases. Given the large number of possible operators it is virtually impossible to isolate the amplitudes and phases of the different operators. (It may be possible to do this in the context of a particular model, in which only a small subset of operators is present.) Fortunately, as we argue below, the strong phases of all NP operators are small relative to those of the SM and can be neglected. As a result, the various NP terms can be combined into a single NP operator, whose amplitude and phase *can* be measured.

To see how this works, consider $B_d^0 \to \phi K_S$. (This is chosen for illustration only – the argument holds for any B decay which receives a significant $\bar{b} \to \bar{s}$ penguin contribution in the SM, and is dominated by a single amplitude.) The SM amplitude for this decay can be written

$$A(B_d^0 \to \phi K_S) = A'_u V_{ub}^* V_{us} + A'_c V_{cb}^* V_{cs} + A'_t V_{tb}^* V_{ts} \ . \tag{5}$$

Here, A'_t arises due to the gluonic penguin amplitude with a t-quark in the loop. Although A'_u and A'_c also receive (small) contributions from the gluonic penguin, they arise mainly as a result of QCD rescattering from the tree operators $\bar{b} \to \bar{s} u \bar{u}$ and $\bar{b} \to \bar{s} c \bar{c}$. The Wilson coefficients for the various contributions imply that $A'_u, A'_c \lesssim 0.5 A'_t$. Note that the size of the rescattered penguin amplitudes is only about 5–10% of that of the tree amplitude. Using CKM unitarity, the amplitude for $B_d^0 \to \phi K_S$ can be written

$$A(B_d^0 \to \phi K_S) = \mathcal{A}'_{ut} \ e^{i\gamma} e^{i\delta'_{ut}} + \mathcal{A}'_{ct} \ e^{i\delta'_{ct}} \approx \mathcal{A}'_{ct} \ e^{i\delta'_{ct}} \ , \tag{6}$$

where $\mathcal{A}'_{ut} \equiv |(A'_u - A'_t) V_{ub}^* V_{us}|$ and $\mathcal{A}'_{ct} \equiv |(A'_c - A'_t) V_{cb}^* V_{cs}|$. The final (approximate) equality arises from the fact that $|V_{ub}^* V_{us} / V_{cb}^* V_{cs}| \simeq 2\%$, so that $\mathcal{A}'_{ut} \ll \mathcal{A}'_{ct}$. The quantity δ'_{ct} is a strong phase; the weak phase is approximately zero.

The principal NP contribution to $B_d^0 \to \phi K_S$ comes from $\bar{s} b \bar{s} s$ (both Lorentz and colour factors are once again suppressed). However, other NP operators, such as $\bar{s} b \, \bar{c} c$, can also contribute to $B_d^0 \to \phi K_S$ through rescattering. The full amplitude for this decay can therefore be written

$$\begin{gathered} A_{\phi K_S} = \mathcal{A}'_{ct} \ e^{i\delta'_{ct}} + \mathcal{A}^{dir}_{NP} + \mathcal{A}^{rescatt}_{NP} \ , \\ \mathcal{A}^{dir}_{NP} \equiv \sum_i A_i e^{i\phi_i^{ss}} e^{i\delta_i} \quad , \qquad \mathcal{A}^{rescatt}_{NP} \equiv \sum_i \epsilon_i B_i e^{i\xi_i} e^{i\sigma_i} \ . \end{gathered} \tag{7}$$

In the above, $\mathcal{A}^{dir}_{NP}$ is the contribution from all NP operators of the form $\bar{s} \Gamma_i b \, \bar{s} \Gamma_j s$ ($\Gamma_{i,j}$ represent Lorentz structures, and colour indices are suppressed), while $\mathcal{A}^{rescatt}_{NP}$ is the contribution from all NP operators of the form $\bar{s} \Gamma_i b \, \bar{q} \Gamma_j q$ ($q \neq s$). In the latter case, the decays $\bar{b} \to \bar{s} q \bar{q}$ ($q \neq s$) contribute to $\bar{b} \to \bar{s} s \bar{s}$ through rescattering. Similarly, $\mathcal{A}^{dir}_{NP}$ includes the "self-rescattering" contributions of $\bar{b} \to \bar{s} s \bar{s}$ to $\bar{b} \to \bar{s} s \bar{s}$. The NP weak phases are ϕ_i^s and ξ_i, while δ_i and σ_i are the NP strong phases.

At this point, it is useful to discuss rescattering in somewhat more detail. As noted above, in the SM, for decays described by $\bar{b} \to \bar{s}$ transitions, the rescattering comes mainly from the tree-level decay $\bar{b} \to \bar{s} c \bar{c}$. Although the rescattered "penguin" amplitudes A'_u and A'_c are only about 5–10% as large as the amplitude which causes the rescattering, they are still of the same order as A'_t [see Eq. (5)]. That is, the SM rescattering effects are not small. In particular, since it is rescattering which is the principal source of strong phases, the phase δ'_{ct} in Eq. (7) can be sizeable.

Now, the new-physics rescattering arises from the NP operators. As in the SM, the rescattered amplitude is suppressed by $\epsilon_i \sim$ 5–10% relative to the operator causing the recattering. Thus, although $B_i \sim A_i$ in Eq. (7), $|\mathcal{A}^{rescatt}_{NP}|$ is only 5–10%

as large as $|\mathcal{A}^{dir}_{NP}|$. (The rescattered contributions in $\mathcal{A}^{dir}_{NP}$ are similarly suppressed.) However, the NP operators are assumed to be of the same size as the SM $\bar{b} \to \bar{s}$ penguin amplitude $\mathcal{A}'_{ct}$ [Eq. (7)]. Therefore $\mathcal{A}^{rescatt}_{NP}$ is negligible compared to $\mathcal{A}'_{ct}$ and $\mathcal{A}^{dir}_{NP}$. In addition, we note that the NP strong phase δ_i in $\mathcal{A}^{dir}_{NP}$ vanishes in the limit of no rescattering. Since, as we have argued, this NP rescattering is small, we have $\delta_i \ll \delta'_{ct}$, i.e. the NP strong phases are negligible compared to those of the SM.

These approximations lead to a considerably simpler structure for Eq. (7):

$$\begin{aligned} A_{\phi K_S} &\approx \mathcal{A}'_{ct}\, e^{i\delta'_{ct}} + \mathcal{A}^{dir}_{NP} ~, \\ \mathcal{A}^{dir}_{NP} &\equiv \sum_i A_i e^{i\phi^s_i} = \mathcal{A}^s_{NP} e^{i\Phi_s} ~, \end{aligned} \tag{8}$$

where we have summed up the new physics contributions into a single amplitude. The important point here is that all the NP weak phases come only from operators of the type $O_{\bar{s}s} = \bar{s}\Gamma_i b\, \bar{s}\Gamma_j s$, and so the effective weak phase carries the subscript "s": Φ_s. From Eq. (8) we have

$$\tan\Phi_s = \frac{\sum_i A_i \sin\phi^s_i}{\sum_i A_i \cos\phi^s_i} ~. \tag{9}$$

The above argument holds for the case where there are new-physics contributions to $\bar{b} \to \bar{s}q\bar{q}$ $(q = d, s, c)$. However, $\bar{b} \to \bar{s}u\bar{u}$ is slightly different because the SM decay is not dominated by a single amplitude – there are both tree and penguin contributions. Nevertheless, it is straightforward to show that the above logic still holds: the rescattering in the NP amplitudes to $\bar{b} \to \bar{s}u\bar{u}$ is negligible, so that the NP contributions can be parametrized by a single amplitude $\mathcal{A}^u_{NP}$ and weak phase Φ_u.

Thus, under the assumption that new-physics rescattering is negligible compared to that of the SM, the effects of the NP operators $\bar{s}b\bar{q}q$ can be parametrized in terms of the effective NP amplitudes $\mathcal{A}^q_{NP}$ $(q = u, d, s, c)$ and the corresponding weak phases Φ_q. In the rest of the paper we will show how these NP parameters can be measured.

Note that there may be a possible loophole in the above argument. In the SM, the exchange and annihilation contributions are expected to be quite small, for both $\bar{b} \to \bar{d}$ and $\bar{b} \to \bar{s}$ transitions. However, in some approaches to hadronic B decays, such amplitudes may be chirally enhanced if there are pseudoscalars in the final state [11, 12], with resulting large strong phases. Hence annihilation-type topologies generated by NP operators may also lead to large strong phases. On the other hand, such chiral enhancements are not present for vector-vector final states and so the above arguments regarding small NP strong phases are applicable here. Ultimately, the size of exchange and annihilation diagrams is an experimental question, and can be tested by the measurement of decays such as $B^0_d \to D^+_s D^-_s$ and $B^0_d \to K^+K^-$.

In general, we take the effective new-physics phases Φ_q to be flavour non-universal. That is, we assume that the phases for different underlying quark transitions, $\bar{b} \to$

$\bar{s}q\bar{q}$, are not related. This occurs in many models of NP, such as supersymmetry with R-parity-violating terms [13]. However, there are also NP models, such as those including a flavour-changing Z or Z' coupling [14], in which the phases *are* related. This shows that the measurement of the Φ_q will be very useful in identifying the new physics, or at least excluding certain NP models. Note that if *all* the NP operators have the same weak phase φ, one has $\Phi_q = \varphi$, and this phase is process-universal as well as flavour-universal. In this case one can simplify Eq. 7 as

$$\begin{aligned} A_{\phi K_S} &= \mathcal{A}'_{ct}\ e^{i\delta'_{ct}} + \mathcal{A}^s_{NP}\ e^{i\delta_{NP}}\ e^{i\varphi}, \\ \mathcal{A}^s_{NP}\ e^{i\delta_{NP}} &= \textstyle\sum_i A_i e^{i\delta_i} + \sum_i \epsilon_i B_i e^{i\sigma_i}\ . \end{aligned} \tag{10}$$

Factoring out the strong phase δ_{NP}, it is clear that we can cast Eq. (10) in the same form as Eq. (8) without any dynamical input about NP strong phases.

Above we showed that the new-physics effects can be parametrized in terms of a few effective NP parameters. We now describe a method for *measuring* these parameters. This technique closely resembles that of Ref. [15], which we recently proposed for extracting CP phase information. Here we turn this method around. As above, we assume that NP is present only in decays with large $\bar{b} \to \bar{s}$ penguin amplitudes. We further assume that the SM CP phase information is known: these phases can be measured using processes which do not involve large $\bar{b} \to \bar{s}$ penguin amplitudes. In this case, the method can be used to extract the NP parameters.

In order to illustrate the method, we consider a specific pair of B decays. It is straightforward to adapt the technique to other processes. Consider $B_s^0 \to K^0 \bar{K}^0$. In the SM, this decay is dominated by a single $\bar{b} \to \bar{s} d\bar{d}$ penguin decay amplitude. Including new physics, the amplitude for $B_s^0 \to K^0 \bar{K}^0$ can be written as [see Eq. (8)]

$$A(B_s^0 \to K^0 \bar{K}^0) \equiv A = \mathcal{A}'_{ct}\ e^{i\delta'_{ct}} + \mathcal{A}^d_{NP} e^{i\Phi_d}\ , \tag{11}$$

where $\mathcal{A}'_{ct}$ and $\mathcal{A}^d_{NP}$ are the SM and NP amplitudes, respectively. Similarly, δ'_{ct} and Φ_d are the SM strong phase and NP weak phase, respectively. The NP phase is defined analogously to Eq. (9). The amplitude for the CP-conjugate process, $\overline{A}$, can be obtained from the above by changing the sign of Φ_d.

Since the final state $K^0 \bar{K}^0$ is accessible to both B_s^0 and $\overline{B_s^0}$ mesons, one can consider indirect (mixing-induced) CP violation. The time-dependent measurement of $B_s^0(t) \to K^0 \bar{K}^0$ allows one to obtain the three observables

$$\begin{aligned} B &\equiv \frac{1}{2}\left(|A|^2 + |\overline{A}|^2\right) = (\mathcal{A}'_{ct})^2 + (\mathcal{A}^d_{NP})^2 + 2\mathcal{A}'_{ct}\,\mathcal{A}^d_{NP} \cos\delta'_{ct} \cos\Phi_d\ , \\ a_{dir} &\equiv \frac{1}{2}\left(|A|^2 - |\overline{A}|^2\right) = 2\mathcal{A}'_{ct}\,\mathcal{A}^d_{NP} \sin\delta'_{ct} \sin\Phi_d\ , \\ a_I &\equiv \mathrm{Im}\left(e^{-2i\phi_{B_s}} A^* \overline{A}\right) = -(\mathcal{A}'_{ct})^2 \sin 2\phi_{B_s} - 2\mathcal{A}'_{ct}\,\mathcal{A}^d_{NP} \cos\delta'_{ct} \sin(2\phi_{B_s} + \Phi_d) \\ &\qquad\qquad\qquad\qquad - (\mathcal{A}^d_{NP})^2 \sin(2\phi_{B_s} + 2\Phi_d)\ . \end{aligned} \tag{12}$$

It is useful to define a fourth observable:

$$
\begin{aligned}
a_{R} \;\equiv\; \mathrm{Re}\left(e^{-2i\phi_{B_s}} A^* \overline{A}\right) &= (\mathcal{A}'_{ct})^2 \cos 2\phi_{B_s} + 2\mathcal{A}'_{ct}\, \mathcal{A}^{d}_{NP} \cos\delta'_{ct} \cos(2\phi_{B_s} + \Phi_d) \\
&\quad + (\mathcal{A}^{d}_{NP})^2 \cos(2\phi_{B_s} + 2\Phi_d) \ . \qquad (13)
\end{aligned}
$$

The quantity a_R is not independent of the other three observables:

$$
{a_R}^2 = B^2 - {a_{dir}}^2 - {a_I}^2 \ . \qquad (14)
$$

Thus, one can obtain a_R from measurements of B, a_{dir} and a_I, up to a sign ambiguity.

In the above, ϕ_{B_s} is the phase of B_s^0–$\overline{B_s^0}$ mixing. In general, NP which affects $\bar{b} \to \bar{s}$ transitions will also contribute to B_s^0–$\overline{B_s^0}$ mixing, i.e. one will have NP operators of the form $\bar{s}b\,\bar{b}s$. In this case, the phase of B_s^0–$\overline{B_s^0}$ mixing may well differ from its SM value ($\simeq 0$) due to the presence of NP. The standard way to measure this mixing phase is through CP violation in $B_s^0(t) \to J/\psi\eta$ (or $B_s^0(t) \to J/\psi\phi$). However, there is a potential problem here: this decay receives NP contributions from $O^c_{NP} \sim \bar{s}b\bar{c}c$ operators (as usual, the Lorentz and colour structures have been suppressed), so that there may be effects from these NP operators in any process involving B_s^0–$\overline{B_s^0}$ mixing.

The solution to this problem can be found by considering B_d^0–$\overline{B_d^0}$ mixing. The phase of this mixing is unaffected by new physics and thus takes its SM value, β. The canonical way to measure this angle is via CP violation in $B_d^0(t) \to J/\psi K_S$. However, this decay also receives NP contributions from O^c_{NP} operators. On the other hand, the value of β extracted from $B_d^0(t) \to J/\psi K_S$ is in line with SM expectations. This strongly suggests that any O^c_{NP} contributions to this decay are quite small. Now, the non-strange part of the η wavefunction has a negligible contribution to $\langle J/\psi\eta |\, O^c_{NP} \,| B_s^0\rangle$. Thus, this matrix element can be related by flavour SU(3) to $\langle J/\psi K_S |\, O^c_{NP} \,| B_d^0\rangle$ (up to a mixing angle). That is, both matrix elements are very small. In other words, we do not expect significant O^c_{NP} contributions to $B_s^0(t) \to J/\psi\eta$, and the phase of B_s^0–$\overline{B_s^0}$ mixing can be measured through CP violation in this decay, even in the presence of NP.

We have already noted that there are many signals of new physics in B decays. Indeed, the expressions for a_{dir} and a_I in Eq. (12) give us clear signals of NP. Since $B_s^0 \to K^0\bar{K}^0$ is dominated by a single decay in the SM, the direct CP asymmetry is predicted to vanish. Furthermore, the indirect CP asymmetry is expected to measure the mixing phase $\phi_{B_s} \simeq 0$. Thus, if it is found that $a_{dir} \neq 0$, or that ϕ_{B_s} does not take its SM value, this would be a smoking-gun signal of NP. Note also that, if it happens that the SM strong phases are small, a_{dir} may be unmeasurable. In this case, a better signal of new physics is the measurement of T-violating triple-product correlations in the corresponding vector-vector final states [6]. This brief discussion illustrates that there are indeed many ways of detecting the presence of NP. However, it must also be stressed that these signals do not, by themselves, allow the measurement of the NP parameters.

The three independent observables of Eqs. (12) and (13) depend on four unknown theoretical parameters: $\mathcal{A}^d_{NP}$, $\mathcal{A}'_{ct}$, δ'_{ct} and Φ_d. Therefore one cannot obtain information about the new-physics parameters $\mathcal{A}^d_{NP}$ and Φ_d from these measurements. However, one can partially solve the equations to obtain

$$(\mathcal{A}'_{ct})^2 = \frac{a_R \cos(2\phi_{B_s} + 2\Phi_d) - a_I \sin(2\phi_{B_s} + 2\Phi_d) - B}{\cos 2\Phi_d - 1} \ . \tag{15}$$

Thus, if we knew $\mathcal{A}'_{ct}$, we could solve for Φ_d.

In order to get $\mathcal{A}'_{ct}$ we consider the partner process $B^0_d \to K^0 \bar{K}^0$, involving a $\bar{b} \to \bar{d}$ penguin amplitude. In the SM this decay is related by SU(3) symmetry to $B^0_s \to K^0 \bar{K}^0$ [16]. Since $\bar{b} \to \bar{s}$ transitions are not involved, the amplitude for $B^0_d \to K^0 \bar{K}^0$ receives only SM contributions, and is given by

$$\begin{aligned} A(B^0_d \to K^0 \bar{K}^0) &= A_u V^*_{ub} V_{ud} + A_c V^*_{cb} V_{cd} + A_t V^*_{tb} V_{td} \\ &= (A_u - A_t) V^*_{ub} V_{ud} + (A_c - A_t) V^*_{cb} V_{cd} \\ &\equiv \mathcal{A}_{ut} \, e^{i\gamma} e^{i\delta_{ut}} + \mathcal{A}_{ct} \, e^{i\delta_{ct}} \ , \end{aligned} \tag{16}$$

where $\mathcal{A}_{ut} \equiv |(A_u - A_t) V^*_{ub} V_{ud}|$, $\mathcal{A}_{ct} \equiv |(A_c - A_t) V^*_{cb} V_{cd}|$, and we have explicitly written the strong phases δ_{ut} and δ_{ct}, as well as the weak phase γ.

As with $B^0_s \to K^0 \bar{K}^0$, the time-dependent measurement of $B^0_d(t) \to K^0 \bar{K}^0$ allows one to obtain three independent observables [Eqs. (12) and (13)]. These observables depend on five theoretical quantities: $\mathcal{A}_{ct}$, $\mathcal{A}_{ut}$, $\delta \equiv \delta_{ut} - \delta_{ct}$, γ and the mixing phase ϕ_{B_d}. However, as discussed above, ϕ_{B_d} can be measured independently using $B^0_d(t) \to J/\psi K_S$. The weak phase γ can also be measured in B decays which are unaffected by new physics in $\bar{b} \to \bar{s}$ penguin amplitudes. For example, it can be obtained from $B^\pm \to DK$ decays [17]. Alternatively, the angle α can be extracted from $B \to \pi\pi$ [18] or $B \to \rho\pi$ decays [19], and γ can be obtained using $\gamma = \pi - \beta - \gamma$. Given that these CP phases can be measured independently, the three observables of $B^0_d(t) \to K^0 \bar{K}^0$ now depend on three unknown theoretical parameters, so that the system of equations can be solved.

In particular, one can obtain $\mathcal{A}_{ct}$:

$$\mathcal{A}^2_{ct} = \frac{a_R \cos(2\phi_{B_d} + 2\gamma) - a_I \sin(2\phi_{B^0_d} + 2\gamma) - B}{\cos 2\gamma - 1} \ , \tag{17}$$

where a_R, a_I and B are the observables found in $B^0_d(t) \to K^0 \bar{K}^0$.

The key point is that, in the SU(3) limit, one has

$$\mathcal{A}_{ct} = \lambda \mathcal{A}'_{ct} \ , \tag{18}$$

where $\lambda = 0.22$ is the Cabibbo angle. Thus, using the above relation, the measurement of $B^0_d(t) \to K^0 \bar{K}^0$ gives us $\mathcal{A}'_{ct}$, in which case Eq. (15) can be used to solve for the new physics phase Φ_d. The NP amplitude $\mathcal{A}^d_{NP}$ can also be obtained. There

NP Parameters	$\bar{b} \to \bar{s}$ Decay	$\bar{b} \to \bar{d}$ Decay
Φ_{cc}, $\mathcal{A}^c_{NP}$	$B^0_s(t) \to D^+_s D^-_s$	$B^0_d(t) \to D^+ D^-$
Φ_s, $\mathcal{A}^s_{NP}$	$B^0_d(t) \to \phi K^{*0}$ $B^0_s(t) \to \phi\phi$	$B^0_s(t) \to \phi \bar{K}^{*0}$ $B^0_s(t) \to \phi \bar{K}^{*0}$
Φ_d, $\mathcal{A}^d_{NP}$	$B^0_s(t) \to K^0 \bar{K}^0$ $B^0_s(t) \to K^0 \bar{K}^0$ $B^0_d(t) \to K^{*0} \rho^0$ $B^0_d(t) \to K^{*0} \rho^0$	$B^0_d(t) \to \pi^+ \pi^-$ $B^0_d(t) \to K^0 \bar{K}^0$ $B^0_d(t) \to \rho^0 \rho^0$ $B^0_s(t) \to \bar{K}^{*0} \rho^0$
Φ_u, $\mathcal{A}^u_{NP}$	$B^0_s(t) \to K^+ K^-$	$B^0_d(t) \to \pi^+ \pi^-$

Table 1: The $\bar{b} \to \bar{s}$ B decays and their $\bar{b} \to \bar{d}$ partner processes which can be used to measure the new-physics parameters $\mathcal{A}^q_{NP}$ and Φ_q.

is a theoretical error in Eq. (18) due to SU(3)-breaking effects. However, various methods were discussed in Ref. [15] to reduce this SU(3) breaking. All of these methods are applicable here. In the end, for this particular pair of processes, the theoretical error is estimated to be in the range 5–10%.

Above, we have shown how measurements of the decays $B^0_{d,s}(t) \to K^0 \bar{K}^0$ can be used to measure the NP parameters $\mathcal{A}^d_{NP}$ and Φ_d. The general idea is to use a $\bar{b} \to \bar{s}$ decay which is dominated in the SM by a single decay amplitude, along with its $\bar{b} \to \bar{d}$ partner process. This method can be adapted to other pairs of B decays to measure different NP parameters. (Or one can find alternative ways of measuring $\mathcal{A}^d_{NP}$ and Φ_d.) By choosing the two decays carefully, the theoretical error can be reduced to the level of 5–15%.

Note that it is only quark-level decays $\bar{b} \to \bar{s} q \bar{q}$ ($q = d, s, c$) which are dominated by a single decay amplitude in the SM. However, one can also apply this technique to $\bar{b} \to \bar{s} u \bar{u}$ decays, for which the $\bar{b} \to \bar{s}$ decay receives both tree and penguin contributions in the SM. For example, one can use the pair of decays $B^0_s(t) \to K^+ K^-$ and $B^0_d(t) \to \pi^+ \pi^-$ to extract the NP parameters $\mathcal{A}^u_{NP}$ and Φ_u. However, in this case the theoretical error is considerably larger since one has to make three SU(3) assumptions of the type in Eq. (18).

In Table 1, we present the list of all B decay pairs to which this method can be applied, along with the NP parameters measured. From this Table, we see that all NP parameters can be obtained. A more detailed analysis of these decays is presented in Ref. [20]. Note that only one decay pair in Table 1 involves only B^0_d decays. The others will require the time-dependent measurement of B^0_s decays. However, this may be difficult experimentally, as B^0_s–$\overline{B^0_s}$ mixing is large. For this reason the decay pair $B^0_d(t) \to K^{*0} \rho^0$ and $B^0_d(t) \to \rho^0 \rho^0$ may be the most promising for measuring NP parameters.

Table 1 lists 8 pairs of B decays. In fact, there are more decay pairs, since

many of the particles in the final states can be observed as either pseudoscalar (P) or vector (V) mesons. Note that certain decays are written in terms of VV final states, while others are have PP states. There are three reasons for this. First, some decays involve a final-state π^0. However, experimentally it will be necessary to find the decay vertices of the final particles. This is virtually impossible for a π^0, and so we always use a ρ^0. Second, some pairs of decays are related by SU(3) in the SM only if an $(s\bar{s})$ quark pair is used. However, there are no P's which are pure $(s\bar{s})$. The mesons η and η' have an $(s\bar{s})$ component, but they also have significant $(u\bar{u})$ and $(d\bar{d})$ pieces. As a result the $\bar{b} \to \bar{s}$ and $\bar{b} \to \bar{d}$ decays are not really related by SU(3) in the SM if the final state involves an η or η'. We therefore consider instead the vector meson ϕ which is essentially a pure $(s\bar{s})$ quark state. Finally, we require that both B^0 and $\bar{B}^0$ be able to decay to the final state. This cannot happen if the final state contains a single K^0 (or $\bar{K}^0$) meson. However, it can occur if this final-state particle is an excited neutral kaon. In this case one decay involves K^{*0}, while the other has $\bar{K}^{*0}$. Assuming that the vector meson is detected via its decay to $K_s\pi^0$ (as in the measurement of $\sin 2\beta$ via $B_d^0(t) \to J/\psi K^*$), then both B^0 and $\bar{B}^0$ can decay to the same final state.

Apart from these three restrictions, the final-state particles can be taken to be either pseudoscalar or vector. Indeed, it will be useful to measure the NP parameters in modes with PP, PV and VV final-state particles, since different NP operators are probed in these decays. For example, within factorization, certain scalar operators in Eq. 4 (i.e. those whose coefficients are $f^{AB}_{q,(1,2)}$) cannot contribute to PV or VV states if their amplitudes involve the matrix element $\langle V|\, \bar{q}\gamma_{L,R}q\, |0\rangle$. In general, the matrix element of a given operator will be different for the various PP, PV and VV final states. Thus, the measurement of the NP parameters in different modes will provide some clues as to which NP operators are present.

Note also that, in general, the value of Φ_q extracted from two distinct decay pairs with the same underlying $\bar{b} \to \bar{s}q\bar{q}$ transition will be different. There are two reasons for this. First, certain operators which contribute to one process may not contribute in the same form in another. (For example, one decay might be colour-suppressed, while the other is colour-allowed.) Second, in general, the matrix elements A_i of the various operators depend on the final states considered. Thus, the value of the NP phase Φ_q depends on the particular decay pair used. However, if all NP operators for the quark-level process $\bar{b} \to \bar{s}q\bar{q}$ have the same weak phase ϕ^q, then the NP phase Φ_q will be the same for all decays governed by the same quark-level process. Hence it is important to measure the phase Φ_q in more than one pair of processes with the same underlying quark transition. If the effective phases are different then it would be a clear signal of more than one NP amplitude, with different weak phases, in $\bar{b} \to \bar{s}q\bar{q}$.

It is also important to measure the NP phases Φ_q for each of $q = u, d, s, c$. As noted earlier, in some NP models, the phases for the different underlying quark transitions $\bar{b} \to \bar{s}q\bar{q}$ are related, so that the NP phase is independent of the quark

flavour. The measurement of the Φ_q would thus allow us to distinguish between NP models.

In summary, it is well known that there are many signals of new physics (NP) which can be found by measuring CP violation in the B system. However, it is usually assumed that one cannot identify the NP – this will have to wait for high-energy colliders which can produce the new particles directly. In this paper we have shown that this is not completely true. We have presented a technique which allows the *measurement* of NP parameters.

In line with hints from present data, we assume that the new physics contributes only to decays with large $\bar{b} \to \bar{s}$ penguin amplitudes, while decays involving $\bar{b} \to \bar{d}$ penguins are not affected. The NP rescattering effects are shown to be small compared to those of the SM and are neglected. This allows us to greatly simplify the form of the NP contributions. In particular, independent of the type of underlying NP, we can parametrize all NP effects in terms of effective NP amplitudes $\mathcal{A}^q_{NP}$ and weak phases Φ_q ($q = u, d, s, c$).

We have shown that one can obtain each of the $\mathcal{A}^q_{NP}$ and Φ_q by using measurements of pairs of B decays. One decay has a large $\bar{b} \to \bar{s}$ penguin component and is (usually) dominated by a single amplitude. It receives a new-physics contribution. The partner process has a $\bar{b} \to \bar{d}$ penguin contribution and is related to the first decay by flavour SU(3) in the SM. It is unaffected by NP. Assuming that the SM CP phases are known independently, the measurements of these two B decays allow one to extract $\mathcal{A}^q_{NP}$ and Φ_q. The theoretical error due to SU(3) breaking can be reduced to the level of 5–15% for $q = d, s, c$, but is larger for $q = u$.

In general, different NP models lead to different patterns of the NP parameters $\mathcal{A}^q_{NP}$ and Φ_q. Thus, the measurement of the NP parameters can rule out certain models and point towards others. We will therefore have a partial identification of the NP, before measurements at high-energy colliders.

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